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Can $\gamma\gamma \rightarrow Z_L Z_L$ serve as a probe of the electroweak symmetry breaking sector ?

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Abstract

We investigate the feasibility of distinguishing among different models of electroweak symmetry breaking by studying the process $\gamma\gamma \rightarrow Z_L Z_L$ at photon colliders. For models with a low mass Higgs-like scalar resonance, the s -channel contribution provides a distinct resonance structure and large cross sections. However, the absence of a resonance structure in the cross section for the above process does not discriminate in principle between the existence of either a heavy Higgs boson or a vector resonance of a non-linearly realized symmetry.

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1 Introduction

The Standard Model has been tested to a remarkable accuracy at LEP but there is still very little to be said about its electroweak symmetry breaking (EWSB) sector apart from a lower limit on the Higgs boson mass, $M_H > 63.5$ GeV [1]. On the theoretical side, the existence of a Higgs boson by itself makes it difficult to understand the fine-tuning that protects its mass to be as large as the Planck mass. There are two possible ways to avoid this so-called hierarchy problem [2]. One could either have a supersymmetric theory, in which quantum corrections to the Higgs mass are only logarithmic instead of quadratic, or one could have a strongly interacting underlying theory which has the EWSB sector as its low-energy effective theory. In spite of some new indirect evidence in favor of supersymmetric models [3], here we will concentrate on the second approach, which has rich experimental consequences at the TeV scale that could be tested at the SSC and LHC.

The presence of this strongly interacting sector would manifest itself primarily in the scattering of the longitudinal components of the weak gauge bosons W_L, Z_L , which at high energies behave as the pseudo Nambu-Goldstone bosons originating from the global symmetry breaking occurring in the EWSB sector, $SU(2)_L \times SU(2)_R \rightarrow SU(2)_V$. The residual $SU(2)_V$ symmetry is responsible for keeping $\rho = M_W/(M_Z \cos \theta) = 1$ at tree level. Therefore, W_L 's and Z_L 's behave as techni-pions of this underlying strong dynamics and we can describe their interactions by using chiral lagrangian techniques. If this scenario is correct, it is very plausible that resonances would also play a role in describing physical processes. In fact, amplitudes derived solely from the pion sector of a chiral lagrangian grow with energy and violate unitarity rather quickly; the presence of resonances tend to unitarize these amplitudes. At this point, there are two types of resonances that could appear : scalar and vector resonances. The usual Higgs sector is an example of the former choice; however, in QCD we know that vector resonances are more important in restoring unitarity, *i.e.*, in saturating scattering amplitudes. Therefore, we should keep an open mind with respect to what type of resonance would show up in $W_L W_L$ scattering. Both types of resonances can be nicely described in terms of chiral lagrangians. A detailed study of consequences of chiral lagrangians with either scalar or vector resonances (and of other models as well) in $W_L W_L$ scattering was performed recently by J. Bagger *et al.* [4].

In this paper we explore the possibility of distinguishing between chiral lagrangians with scalar or vector resonances by studying their contribution to the process $\gamma\gamma \rightarrow Z_L Z_L$. The suggestion of obtaining high energy γ beams from backscattered laser and the smallness of the above process in the Standard Model makes it in principle a good window to explore physics beyond the Standard Model. There has been some

recent activity in the study of $\gamma\gamma \rightarrow Z_L Z_L$. Abbasabadi et al. [5] used a K-matrix unitarized amplitude to study this process in the Standard Model, i.e., a model with a scalar resonance. They found a substantial correction to the same calculation without unitarization [6] only for $M_H > 5$ TeV, in accord to the fact that the Standard Model amplitudes violate unitarity for $M_H > \mathcal{O}(1\text{TeV})$. This process was also considered in the context of chiral lagrangians without any resonances in ref. [7] and in the context of anomalous gauge couplings in ref. [8].

In the next section we compute the amplitudes for $\gamma\gamma \rightarrow Z_L Z_L$ in chiral lagrangian models with no resonance, a scalar resonance and a vector resonance. In section 3 we compare the cross sections arising from these different models. In section 4 we examine the possibility of distinguishing these models in a realistic $\gamma\gamma$ machine and we conclude in section 5.

2 Models and amplitudes

2.1 Chiral lagrangian without resonances

This model describes the interaction of the pseudo Nambu-Goldstone bosons w^a ($a = 1, 2, 3$) among themselves and the photon field A_μ given by the lagrangian

$$\mathcal{L} = \frac{v^2}{4} \text{Tr}[D_\mu U D^\mu U^\dagger] - \frac{1}{4} (F_{\mu\nu})^2, \quad (1)$$

where

$$U = \exp[2iw/v] \quad , \quad w = w^a \tau^a / 2 \quad , \quad v = 246 \text{ GeV}, \quad (2)$$

and the covariant derivative is given by :

$$D_\mu U = \partial_\mu U + ieA_\mu [Q, U]. \quad (3)$$

A few words must be said about the choice for the charge matrix Q . The chiral lagrangian possesses the symmetries of the underlying theory involving fermions. In the general case where these symmetries are local $G_L \times G_R$ and the fundamental fermions transform as

$$\psi_L \rightarrow L\psi_L \quad , \quad \psi_R \rightarrow R\psi_R \quad (4)$$

where $L \in G_L$ and $R \in G_R$, the matrix U transforms as :

$$U \rightarrow LUR^\dagger. \quad (5)$$

Introducing the gauge fields l_μ and r_μ transforming as

$$l_\mu \rightarrow Ll_\mu L^\dagger + i(\partial_\mu L)L^\dagger \quad , \quad r_\mu \rightarrow Rr_\mu R^\dagger + i(\partial_\mu R)R^\dagger \quad (6)$$

the covariant derivative becomes :

$$D_\mu U = \partial_\mu U + i l_\mu U - i U r_\mu. \quad (7)$$

In the case of QCD with only up and down quarks, the electromagnetic interactions are introduced by the identifications :

$$l_\mu = r_\mu = e A_\mu Q \quad , \quad Q = \begin{pmatrix} 2/3 & 0 \\ 0 & -1/3 \end{pmatrix}, \quad (8)$$

recovering Eq. (3).

In the case of the Standard Model, we introduce $SU(2)_L \times U(1)_Y$ gauge interactions by choosing :

$$\begin{aligned} l_\mu &= g' \frac{Y_L}{2} B_\mu + g \frac{\tau^a}{2} W_\mu^a \\ r_\mu &= g' \frac{Y_R}{2} B_\mu \end{aligned} \quad (9)$$

which correspond to a charge matrix :

$$Q = \begin{pmatrix} 1/2 & 0 \\ 0 & -1/2 \end{pmatrix} + \frac{Y_L}{2} \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}, \quad (10)$$

This charge matrix should reflect the electroweak quantum numbers of the underlying techni-fermions; it should be really considered as a free parameter of the model. Here we'll parametrize the charge matrix by Y_L , the hypercharge of the techni-fermion doublet. We have checked that all interactions involving only photons and w fields described by the lagrangian Eq.(1) are independent of Y_L . For the general case of $Y_L \neq 0$, extra heavy fermions would be required to cancel undesirable anomalies.

We are interested in the process $\gamma\gamma \rightarrow Z_L Z_L$. Using the equivalence theorem, which was proved to all orders in the weak coupling constant in any R_ξ gauge [9], we have for the scattering amplitude at a given center-of-mass energy \sqrt{s} (denoting $z = w^3$) :

$$A(\gamma\gamma \rightarrow Z_L Z_L) = A(\gamma\gamma \rightarrow zz) + \mathcal{O}(M_Z^2/s). \quad (11)$$

In the following we'll write :

$$A(\gamma\gamma \rightarrow zz) = \epsilon_1^\mu \epsilon_2^\nu T_{\mu\nu}, \quad (12)$$

where $\epsilon_{1,2}$ are the polarization 4-vectors of the initial photons.

The calculation of the process $\gamma\gamma \rightarrow zz$ is straightforward but tedious. From the lagrangian Eq.(1) one obtains the relevant couplings γww , $\gamma\gamma ww$, www , $\gamma\gamma www$

and $\gamma w w w$ ($w = w^\pm, z$) that enter in the computation of the relevant one-loop Feynman diagrams (there are no tree-level contributions in this model). The result for this non-resonance (NR) model is [10, 11] :

$$T_{\mu\nu}^{NR} = -i \frac{\alpha}{2\pi} \left(g_{\mu\nu} - \frac{k_{1\nu} k_{2\mu}}{k_1 \cdot k_2} \right) \frac{s}{v^2}, \quad (13)$$

where k_1 and k_2 are momentum 4-vectors associated with the photons. The result is finite because there is no tree-level contribution to this process at $\mathcal{O}(p^4)$ that would absorb the divergencies. Noticing that in this model,

$$A_{NR}(w^+ w^- \rightarrow z z) = i s / v^2, \quad (14)$$

we can write :

$$T_{\mu\nu}^{NR} = -\frac{\alpha}{2\pi} \left(g_{\mu\nu} - \frac{k_{1\nu} k_{2\mu}}{k_1 \cdot k_2} \right) A_{NR}(w^+ w^- \rightarrow z z), \quad (15)$$

which can be interpreted as a rescattering $\gamma\gamma \rightarrow w^+ w^- \rightarrow z z$.

This simple model should not be trusted at energies $\sqrt{s} > \mathcal{O}(4\pi v)$, where unitarity is violated in some amplitudes for ww scattering and all the terms in the derivative expansion become of the same order. The effects of higher order terms can be estimated by introducing resonances that saturate the scattering amplitudes or by unitarizing these amplitudes in an *ad hoc* manner. The latter approach was used by the authors of ref. [10, 12] in the case of $\gamma\gamma \rightarrow \pi^0 \pi^0$.

The result Eq.(15) will serve as an *ansatz* to study the influence of resonances on the process $\gamma\gamma \rightarrow Z_L Z_L$, *i.e.*, we will assume that this relation is valid in the other models discussed below as well. In the next subsections we'll examine two classes of models : a Higgs-like model with a scalar resonance and a technicolor-like model with a vector resonance.

2.2 Chiral lagrangian with a scalar resonance

Here we consider the most general model of a scalar particle H coupled in a chiral invariant way [13] :

$$\mathcal{L}_S = \frac{v^2}{4} \text{Tr}[D_\mu U D^\mu U^\dagger] - \frac{1}{4} (F_{\mu\nu})^2 + \frac{1}{2} \partial_\mu H \partial^\mu H - \frac{1}{2} M_H^2 H^2 + \frac{1}{2} g_H v H \text{Tr}[D_\mu U D^\mu U^\dagger]. \quad (16)$$

The scalar resonance is parametrized by its mass M_H and by a coupling constant g_H . Its width is given by :

$$\Gamma_H = \frac{3g_H^2 M_H^3}{32\pi v^2} \quad (17)$$

and the relevant scattering amplitude in this model reads :

$$A_H(w^+w^- \rightarrow zz) = i \frac{s}{v^2} \left[1 - \frac{g_H^2 s}{s - M_H^2 + iM_H \Gamma_H} \right]. \quad (18)$$

Notice that H reduces to the Standard Model Higgs boson for $g_H = 1$. In what follows we'll restrict ourselves to the case $g_H = 1$. We don't expect any major differences for other cases.

From Eq.(15) we find the contribution of a generic scalar resonance to $\gamma\gamma \rightarrow zz$:

$$T_{\mu\nu}^H = -i \frac{\alpha s}{2\pi v^2} \left(g_{\mu\nu} - \frac{k_{1\nu} k_{2\mu}}{k_1 \cdot k_2} \right) \left[1 - \frac{g_H^2 s}{s - M_H^2 + iM_H \Gamma_H} \right]. \quad (19)$$

For $g_H = 1$, this agrees with the result of Ref. [6], which was obtained in the context of the Standard Model but keeping only couplings of enhanced electroweak strength, i.e., neglecting couplings of the order $\mathcal{O}(g^2)$ compared to $\mathcal{O}(g^2 M_H^2/M_W^2)$.

2.3 Chiral lagrangian with a vector resonance

We introduce a vector resonance in the chiral lagrangian as a gauge vector boson of a local $SU(2)$ group by means of the so-called hidden symmetry approach [14], which has been successful in describing the properties of vector mesons in QCD. It was shown to be equivalent to other different approaches in Ref. [15].

We parametrize the matrix U by :

$$U = \xi_L^\dagger \xi_R \quad (20)$$

with the following transformations under $G_L \times G_R \times SU(2)$:

$$\begin{aligned} \xi_L &\rightarrow h \xi_L L^\dagger \\ \xi_R &\rightarrow h \xi_R R^\dagger \\ U &\rightarrow L U R^\dagger \end{aligned} \quad (21)$$

where $L \in G_L$, $R \in G_R$ and $h \in SU(2)$. Let's assume for the moment that G_L and G_R are global groups; external gauge fields are easily introduced by gauging the appropriate subgroups.

The vector resonance $V_\mu = V_\mu^a \frac{\tau^a}{2}$ transforms as a $SU(2)$ gauge field :

$$V_\mu \rightarrow h V_\mu h^\dagger + \frac{i}{g''} h \partial_\mu h^\dagger, \quad (22)$$

where g'' is the coupling constant associated with the local $SU(2)$. We define the covariant derivative

$$D_\mu \xi_{L,R} = \partial_\mu \xi_{L,R} - ig'' V_\mu \xi_{L,R} \quad (23)$$

such that

$$D_\mu \xi_L \rightarrow h D_\mu \xi_L L^\dagger \quad , \quad D_\mu \xi_R \rightarrow h D_\mu \xi_R R^\dagger \quad (24)$$

The building blocks to construct an invariant lagrangian are :

$$\alpha_{L\mu} = (D_\mu \xi_L) \xi_L^\dagger \quad , \quad \alpha_{R\mu} = (D_\mu \xi_R) \xi_R^\dagger \quad (25)$$

which transform as

$$\alpha_{L\mu} \rightarrow h \alpha_{L\mu} h^\dagger \quad , \quad \alpha_{R\mu} \rightarrow h \alpha_{R\mu} h^\dagger. \quad (26)$$

Finally , the most general lowest order lagrangian which respects the parity-like symmetry $L \leftrightarrow R$ is given by :

$$\mathcal{L}_V = -\frac{1}{4} Tr[(V_{\mu\nu})^2] - \frac{v^2}{4} Tr[(\alpha_{L\mu} - \alpha_{R\mu})^2] - a \frac{v^2}{4} Tr[(\alpha_{L\mu} + \alpha_{R\mu})^2], \quad (27)$$

where $V_{\mu\nu}$ is the non-abelian field-strength for the vector resonance. There are two free parameters in this lagrangian, namely, g'' and a . It can be shown that in the limit $g'' \rightarrow \infty$, the kinetic term for the vector resonance vanishes and the resonance becomes an auxiliary field which is eliminated by its equation of motion, which in turn reduces the above lagrangian to the usual non-linear σ model lagrangian :

$$\mathcal{L}_V \xrightarrow{g'' \rightarrow \infty} \frac{v^2}{4} Tr[\partial_\mu U \partial^\mu U^\dagger]. \quad (28)$$

The mass and width of the vector resonance in this model can be derived from the lagrangian after using an $SU(2)$ gauge transformation to set $\xi_L^\dagger = \xi_R = \exp[iw/v]$ (unitary gauge):

$$\begin{aligned} M_V^2 &= ag''^2 v^2 \\ \Gamma_V &= \frac{aM_V^3}{192\pi v^2} \end{aligned} \quad (29)$$

and the scattering amplitude of interest to us is given by [16]:

$$-iA_V(w^+ w^- \rightarrow zz) = \frac{s}{4v^2} (4 - 3a) + \frac{aM_V^2}{4v^2} \left[\frac{u - s}{t - M_V^2} + \frac{t - s}{u - M_V^2} \right]. \quad (30)$$

Notice that the above amplitude reduces to Eq.(14) in the limit $M_V^2 \gg t, u$, reproducing the low-energy theorems.

However, the parity-like operation $L \leftrightarrow R$ is not a symmetry of the underlying theory. It corresponds to a symmetry of the theory under $w(\vec{x}, t) \rightarrow -w(\vec{x}, t)$, which

forbids transitions between even and odd numbers of the pseudo-Goldstone bosons w . However, parity conservation implies in the symmetry $w(\vec{x}, t) \rightarrow -w(-\vec{x}, t)$ and it is possible to write down parity-conserving terms in the lagrangian that violate the $L \leftrightarrow R$ symmetry [17]. In QCD, these terms describe processes like $\rho, \omega \rightarrow \pi\gamma$. An analogue term in the model we are studying is uniquely (up to total derivatives) determined to be [18] :

$$\mathcal{L}_{V\gamma w} = \kappa \frac{eg''}{v} \epsilon^{\mu\nu\rho\sigma} V_\mu^a \partial_\nu A_\rho \partial_\sigma w^b \text{Tr}[Q\{T^a, T^b\}], \quad (31)$$

where κ is a constant. This is the first time that the charge matrix Q becomes relevant. In fact, for Q given by Eq.(10) we find that the above lagrangian is proportional to Y_L . It is interesting to recall that in the context of the non-relativistic quark model one has [19] :

$$\frac{A(\rho \rightarrow \pi\gamma)}{A(\omega \rightarrow \pi\gamma)} = \frac{(e_u + e_d)}{(e_u - e_d)} \quad (32)$$

and there would be no $\rho \rightarrow \pi\gamma$ decay if the quark charge matrix were given by Eq.(10) with $Y_L = 0$. Here we'll include the effect of the trace in an effective coupling κ' . In QCD, $\kappa' \approx 0.03$ from radiative vector meson decays.

With the interaction given by Eq.(31) it is straightforward to compute the contribution to $\gamma\gamma \rightarrow zz$ arising from a t and u -channel exchange of a vector resonance [20]:

$$T_{\mu\nu}^{t+u} = i\left(\frac{\kappa' eg''}{v}\right)^2 \left\{ \left[g_{\mu\nu} - \frac{k_{1\nu} k_{2\mu}}{k_1 \cdot k_2} \right] \left(\frac{st/4}{t - M_V^2} + \frac{su/4}{u - M_V^2} \right) - \left[\frac{s}{2} p_{1\mu} p_{2\nu} + \frac{ut}{4} g_{\mu\nu} + \frac{t}{2} k_{2\mu} p_{1\nu} + \frac{u}{2} k_{1\nu} p_{1\mu} \right] \left(\frac{1}{t - M_V^2} + \frac{1}{u - M_V^2} \right) \right\} \quad (33)$$

and finally the total amplitude in this model is given by :

$$T_{\mu\nu}^V = -\frac{\alpha}{2\pi} \left(g_{\mu\nu} - \frac{k_{1\nu} k_{2\mu}}{k_1 \cdot k_2} \right) A_V(w^+ w^- \rightarrow zz) + T_{\mu\nu}^{t+u}. \quad (34)$$

3 Cross section and comparison between models

For amplitudes written as Eq.(12) the cross section is given by :

$$\frac{d\hat{\sigma}}{dt} = \frac{1}{128\pi s^2} |T_{\mu\nu}|^2 \quad (35)$$

For the no-resonance model and for the scalar model we get :

$$\hat{\sigma}_{NR}(s) = \frac{\alpha^2}{256\pi^3} \frac{s}{v^4} \quad (36)$$

$$\hat{\sigma}_H(s) = \frac{\alpha^2}{256\pi^3} \frac{s}{v^4} \left| 1 - \frac{g_H^2 s}{s - M_H^2 + iM_H \Gamma_H} \right|^2. \quad (37)$$

For the vector model we find :

$$\hat{\sigma}_V(s) = \frac{1}{128\pi s^2} \int_{-s}^0 dt \left[2A^2 + \frac{(ut)^2}{8} B^2 \right], \quad (38)$$

where

$$\begin{aligned} A &= -\frac{\alpha}{2\pi} \left[\frac{s}{4v^2}(4-3a) + \frac{aM_V^2}{4v^2} \left(\frac{u-s}{t-M_V^2} + \frac{t-s}{u-M_V^2} \right) \right] \\ &\quad + \left(\frac{k'eg''}{v} \right)^2 \left[\frac{st/4}{t-M_V^2} + \frac{su/4}{u-M_V^2} \right] \\ B &= \left(\frac{k'eg''}{v} \right)^2 \left[\frac{1}{t-M_V^2} + \frac{1}{u-M_V^2} \right] \end{aligned} \quad (39)$$

In Figs. 1a, 1b we plot these cross sections for typical values of the parameters. We immediately see that the model with a vector resonance is characterized by an *absence* of the signal. This can be understood by the fact that, in contrast to the scalar resonance case, there is no s -channel resonance contribution to the rescattering process $w^+w^- \rightarrow zz$ in the vector resonance model.

A few words are now in order with respect to our calculation in the vector model. At the energies we are considering, for κ' close to its QCD value, the contribution from Eq. (31) is very small, the end result being almost identical to taking $\kappa' = 0$. There is a priori no dynamical reason for not having larger values of κ' (κ' is proportional to the techniquark magnetic moment), which typically results in an enhancement of the cross section at high energies.

Also of interest is the result of our calculation to QCD. In that case, we find a cross section for $\gamma\gamma \rightarrow \pi^0\pi^0$ that presents the features of interference between the two contributions and is considerably below the experimental values. However, when the two contributions are taken separately either the unitarized amplitude ($\kappa' = 0$) or the non-unitarized+vector exchange amplitude ($a = 0$) would fairly reproduce the QCD data. We don't expect this crude model to describe the QCD data accurately since we have not include the effects of other resonances like the ω and A_1 [21].

We also emphasize that Eq.(15) is an *assumption* that, although found correct in the scalar model, may not be valid in the vector model. It actually introduces spurious t - and u -channel poles in addition to the physical ones given by the second term in Eq.(34). In fact, it was recently shown by Pennington and Morgan [22] that only for models in which the amplitude for $\pi^+\pi^- \rightarrow \pi^0\pi^0$ is independent of the initial pions momenta does Eq.(15) holds. This is indeed the case for the non-resonance and scalar models. They have also shown by means of an explicit calculation that, in a model that includes both scalar and vector resonances, Eq.(15) does not reproduce the correct threshold behavior for $\gamma\gamma \rightarrow \pi^0\pi^0$. However, the *numerical* difference that exists

between the ansatz of Eq.(15) and the exact result in different kinematic regimes is still unknown. For instance, Dobado and Peláez [12] used a Padé unitarized amplitude for $\pi^+\pi^- \rightarrow \pi^0\pi^0$, which effectively generates a ρ -meson, to compute $\gamma\gamma \rightarrow \pi^0\pi^0$ via Eq.(15) and found a good agreement with their exact unitarized result (and with data¹).

A more complete calculation would involve computing the contribution of the vector resonances in loops, that is, a complete one-loop analysis of this process in the context of the hidden symmetry model [23]. Since this model is supposed to be a good description of QCD, we don't expect any major differences from this more thorough approach.

4 Implications for photon-photon colliders

There has been recent interest in the possibility of producing a high energy photon-photon collider by back-scattering a high intensity laser beam off a high energy electron beam. Most of the scattered photons have their direction close to the original electron beam and carry a large fraction of the initial electron energy. The probability of finding a photon carrying an energy fraction x of the initial electron is given by [24] :

$$F_{\gamma/e}(x) = \frac{N(x, \xi)}{D(\xi)} \quad (40)$$

where ξ is defined in terms of the electron mass m_e , the electron initial energy E and the energy of the photon in the laser beam ω_0 :

$$\xi = \frac{4E\omega_0}{m_e^2} \quad (41)$$

and

$$N(x, \xi) = 1 - x + \frac{1}{1-x} - \frac{4x}{\xi(1-x)} + \frac{4x^2}{\xi^2(1-x)^2} \quad (42)$$

$$\begin{aligned} D(\xi) &= \int_0^{x_m} dx N(x, \xi) \\ &= \left[1 - \frac{4}{\xi} - \frac{8}{\xi^2} \right] \ln(1 + \xi) + \left(\frac{1}{2} + \frac{8}{\xi} - \frac{1}{2(1 + \xi)^2} \right) \end{aligned} \quad (43)$$

where $x_m = \frac{\xi}{1+\xi}$ is the maximum energy fraction carried by the scattered photon.

¹In their model there is no direct $\rho\gamma\pi$ vertex, which can upset the good agreement between their results and experimental data.

The differential cross section with respect to the invariant mass M_{ZZ} of the final state $Z_L Z_L$ pair can be written as :

$$\frac{d\sigma}{dM_{ZZ}}(s) = \frac{2M_{ZZ}}{s} \frac{d\mathcal{L}_{\gamma\gamma}}{d\tau}(\tau) \hat{\sigma}(\tau s) \quad (44)$$

where the differential photon luminosity function $\frac{d\mathcal{L}_{\gamma\gamma}}{d\tau}$ is given by :

$$\frac{d\mathcal{L}_{\gamma\gamma}}{d\tau}(\tau) = \int_{\tau/x_m}^{x_m} \frac{dx}{x} F_{\gamma/e}(x) F_{\gamma/e}(\tau/x), \quad (45)$$

where $\tau = \hat{s}/s = M_{ZZ}^2/s$.

We illustrate our results by showing in Fig. 2 these differential cross sections for a $\gamma\gamma$ collider originated from a $\omega_0 = 1.17\text{eV}$ laser being Compton back-scattered from both e^+ and e^- beams of a 1 TeV linear collider. In this case, $\xi = 9.4$ and $x_m = 0.90$. In this example we'll neglect e^+e^- pair production that may occur for $x_m \geq 0.828$. We concentrate in the region $M_{ZZ} \geq 300$ GeV, where the equivalence theorem as well as the use of massless phase space become good approximations. In this figure we compare the results of the vector resonance model with different parameters with a 1 TeV Higgs model, since for a smaller Higgs mass the signal for the scalar model is much larger than for the vector model. The resonant shape of the curves in this figure is an artifact arising from the convolution of a growing point cross section with a photon luminosity function that falls rapidly near the kinematical limit $(\sqrt{\hat{s}})_{max} = 900$ GeV.

5 Conclusion

The cross section arising from the model with a vector resonance is generally small compared to a model with a scalar resonance of mass $M_H < 800$ GeV, where a pole structure is the dominant feature of the signal. However, the absence of a clear resonance structure would not rule out the scalar model since it may be possible that the scalar mass is above the 1 TeV scale, in which case the cross section is small ($\mathcal{O}(\text{fb})$) in both models. The situation becomes even worse when backgrounds are taken into account [25]. Recently it has been shown that the irreducible background from the production of transversally polarized Z pairs is very large [26], reducing the possibilities of finding even a scalar resonance with mass above 300 GeV. As techniques for the determination of the polarization of gauge bosons are improved and high luminosity ($\mathcal{O}(100 \text{ fb}^{-1} \text{ year}^{-1})$) $\gamma\gamma$ colliders become available [27] it may eventually be possible to extract some information on the symmetry breaking sector from the process $\gamma\gamma \rightarrow Z_L Z_L$. The vector resonance model could in principle be better tested at an $e\gamma$ facility because of the s -channel contribution to the subprocess $\gamma W^{+,-} \rightarrow \gamma W^{+,-}$ that exists in addition to the W boson contribution. Work along these lines is in progress.

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Figure Captions

Figure 1a : Point cross section in femtobarns as a function of the $\gamma\gamma$ center-of-mass energy . Solid line: non-resonance model; dotted line: $M_H = 1000$ GeV; dashed line: $M_H = 800$ GeV; dot-dashed line: $M_H = 600$ GeV.

Figure 1b : Point cross section in femtobarns as a function of the $\gamma\gamma$ center-of-mass energy . Solid line: non-resonance model; dotted line: $M_V = 2000$ GeV, $\Gamma_V = 700$ GeV, $\kappa' = 0.2$ (cross section in this case is scaled down by a factor of 10); dashed line: $M_V = 2000$ GeV, $\Gamma_V = 700$ GeV, $\kappa' = 0.03$; dot-dashed line: $M_V = 1000$ GeV, $\Gamma_V = 300$ GeV, $\kappa' = 0.03$.

Figure 2 : Differential cross sections in fb/GeV as a function of the $Z_L Z_L$ invariant mass. Solid line: $M_V = 2000$ GeV, $\Gamma_V = 700$ GeV, $\kappa' = 0.03$; dotted line: $M_V = 1000$ GeV, $\Gamma_V = 300$ GeV, $\kappa' = 0.03$; dashed line: $M_H = 1000$ GeV; dot-dashed line: $M_V = 2000$ GeV, $\Gamma_V = 700$ GeV, $\kappa' = 0.2$.

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